

Dielectron and Diphoton channels in 2HDM

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Abstract

It was observed recently that [1], any particle in the Standard Model cannot decay to e^+e^- and $\gamma\gamma$ final states with comparably measurable branching ratios. This is also true for most extensions of the Standard Model, with the Randall-Sundrum model as an outstanding exception. In this paper, we show that two-Higgs-Doublet-Models (2HDM) yield another possible exception if certain parameters are properly chosen. In addition, we have checked that this model survives the tests of low energy processes, including the anomalous magnetic moment and electric dipole moment of leptons, lepton-flavor-violating decays $\mu^- \rightarrow e^- \gamma$ and $e^- e^+ e^-$.

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I. INTRODUCTION

It was observed recently that [1], no resonance can decay to e^+e^- and $\gamma\gamma$ with comparably measurable branching ratios, if electrons interact only with the electroweak gauge bosons and the Standard Model (SM) Higgs boson. That is to say, if both dielectron and diphoton channels are observed in experiments such as Tevatron and/or LHC, one gets a smoking-gun evidence of new physics beyond the SM. Furthermore, this observation holds true for most conventional extensions of the SM, with the Randall-Sundrum (RS) [2] model as an exception. Specifically, the ratio $Br(h^{KK} \rightarrow \gamma\gamma)/Br(h^{KK} \rightarrow e^+e^-)$ for the Kaluza-Klein (KK) modes h^{KK} of graviton is predicted to be exactly 2 in the original RS-model. In variant RS-models, this ratio could be different, but generically one expects to observe simultaneously e^+e^- and $\gamma\gamma$ final states with comparable rates. We are thus led to examine carefully whether there are other extensions of the SM which can produce similar signals.

First, let us recapitulate the arguments [1] why e^+e^- and $\gamma\gamma$ channels do not appear simultaneously in the SM. Due to conservation of angular momentum, fermionic resonance can decay to neither e^+e^- nor $\gamma\gamma$ final states. Furthermore, Landau-Yang theorem [3] prohibits a spin-1 resonance decaying into two photons, which follows from general assumptions of Lorentz invariance, gauge invariance, and Bose symmetry of the photons. In principle, a spin-0 resonance can decay to e^+e^- either through the electroweak gauge interaction or the Yukawa interaction. However, the gauge channel is helicity suppressed while the Yukawa channel is suppressed by the ratio of electron mass to the vacuum expectation value (VEV) of the SM Higgs. For a resonance of spin-2 or higher spin, e^+e^- channel is forbidden at tree level simply because the resonance can not couple directly to a virtual spin-1 gauge boson. The arguments hold also true for the case of dimuon and diphoton signals, if the muon mass is much smaller than the energy scale of the resonance.

Thus comparably measurable dielectron and diphoton signals clearly imply that leptons must have direct interaction with either a spin-2 (or higher spin) boson or a new scalar beyond the SM Higgs. The former possibility has been examined in [1] where one gets the KK modes of RS graviton as a concrete example. This is actually not hard to understand, as gravity couples universally to energy-momentum tensor of any fields, including electrons and photons. In this letter, we will instead focus on the latter case, by considering models with new scalar fields that can decay into dielectron and diphoton with comparably measurable

rates. In particular, we will concentrate on extensions of the SM with an extra Higgs doublet, which is usually dubbed as the two Higgs doublet models (2HDM).

In 2HDM, there are eight degrees of freedom in the scalar sector to start with. After spontaneously symmetry breaking, five scalars will remain but two of them are charged. Among the three neutral ones, one resembles the SM Higgs and the additional two could play the role of new resonances. These neutral scalars couple to electron directly through Yukawa interactions. To avoid constraints on the Yukawa couplings due to the light electron mass compared to the electroweak scale, the VEV of the additional Higgs doublet must be (almost) zero. This type of model is usually referred to as type III 2HDM, in which the neutral scalars can decay to e^+e^- at tree level with significant rate. Of course, care should be taken to avoid large FCNC and CP violation effects. On the other hand, the $\gamma\gamma$ decay is loop suppressed in these models and smaller than that in the SM due to the absence of contributions from gauge particles. To make a sizable diphoton decay rate, some mechanism for enhancement is needed for this loop level process. In the letter, we will see that such enhancement is still feasible within current experimental limits. Putting all things together, the extra scalars can be made to decay into both e^+e^- and $\gamma\gamma$ final states with comparably measurable branching ratios. As one would have expected, here one has to choose parameters properly and certain level of fine tuning is necessary.

This paper is organized as follows: In Section II, we will show how to get significantly comparable branching ratios for e^+e^- and $\gamma\gamma$ channels in 2HDM. Phenomenological constraints on these models are presented in Section III, including the anomalous magnetic moment of leptons, electric dipole moment(EDM) of leptons, the lepton flavor violating (LFV) processes $\mu^- \rightarrow e^- \gamma$ and $\mu^- \rightarrow e^- e^+ e^-$. Generically leptonic Yukawa couplings in 2HDM are mostly constrained by the electric dipole moment (EDM) of the electron, instead of by the LFV processes. The results are summarized in section IV.

II. COMPARABLE BRANCHING RATIOS IN TYPE III 2HDM

In Type III 2HDM, one scalar doublet resembles the role of the SM Higgs boson, which will be denoted as ϕ_1 . We denoted the additional doublet as ϕ_2 . The general interaction between ϕ_2 and fermions reads

$$\mathcal{L}_{Yukawa} = -\xi_{ij}^U \bar{Q}_{Li} \tilde{\phi}_2 U_{Rj} - \xi_{ij}^D \bar{Q}_{Li} \phi_2 D_{Rj} - \xi_{ij}^E \bar{l}_{Li} \phi_2 E_{Rj} + h.c. \quad (1)$$

with left (right) handed projection $L(R) = (1 \mp \gamma_5)/2$, \bar{Q}_{Li} (l_{Li}) are left handed quark (lepton) doublets, U_{Rj} , D_{Rj} and E_{Rj} are right handed up quark, down quark and lepton singlets, respectively, with family indices i, j . In general, all Yukawa matrices $\xi_{ij}^{U,D,E}$ are non-diagonal and complex¹. This may result in FCNC and CP violating effects, which will be addressed in the next section. The general Higgs potential [4] which spontaneously breaks $SU(2)_L \times U(1)_Y$ down to $U(1)_{EM}$ in 2HDM is

$$\begin{aligned} V(\phi_1, \phi_2) = & \lambda_1(\phi_1^+ \phi_1 - v_1^2)^2 + \lambda_2(\phi_2^+ \phi_2 - v_2^2)^2 + \lambda_3[(\phi_1^+ \phi_1 - v_1^2) + (\phi_2^+ \phi_2 - v_2^2)]^2 \\ & + \lambda_4[(\phi_1^+ \phi_1)(\phi_2^+ \phi_2) - (\phi_1^+ \phi_2)(\phi_2^+ \phi_1)] + \lambda_5[Re(\phi_1^+ \phi_2) - v_1 v_2 \cos \xi]^2 \\ & + \lambda_6[Im(\phi_1^+ \phi_2) - v_1 v_2 \sin \xi]^2, \end{aligned} \quad (2)$$

where the λ_i 's are all real parameters.

In type I and type II 2HDM, the non-zero VEVs of ϕ_1 and ϕ_2 make contributions to fermion masses proportional to the corresponding Yukawa couplings in (1). The branching ratio of the e^+e^- final state from scalar particles is severely suppressed by the tiny electron mass. In type III 2HDM, those Yukawa couplings are not confined by such constraints, since ϕ_2 has a (almost) zero VEV. One might hope to find viable models with comparably measurable branching ratios for e^+e^- and $\gamma\gamma$ channels.

Only ϕ_1 has a non-zero VEV while ϕ_2 does not in type III 2HDM. That is,

$$v_1 = \langle \phi_1 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v \end{pmatrix} ; \quad v_2 = \langle \phi_2 \rangle = 0 . \quad (3)$$

We can thus choose a gauge such that

$$\phi_1 = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h \end{pmatrix} ; \quad \phi_2 = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2}H^+ \\ H_1 + iH_2 \end{pmatrix} , \quad (4)$$

In this gauge, the potential in (2) can be re-expressed as,

$$\begin{aligned} V(\phi_1, \phi_2) = & \frac{1}{2} [2(\lambda_1 + \lambda_3)v^2] h^2 + \left(\frac{1}{2} \lambda_4 v^2 \right) H^- H^+ + \frac{1}{2} \left(\frac{1}{2} \lambda_5 v^2 \right) H_1^2 \\ & + \frac{1}{2} \left(\frac{1}{2} \lambda_6 v^2 \right) H_2^2 + (\text{interaction terms}), \end{aligned} \quad (5)$$

¹ The Yukawa matrices for interactions between ϕ_1 and fermions are diagonalized here, so all fermions are already in their mass eigenstates.

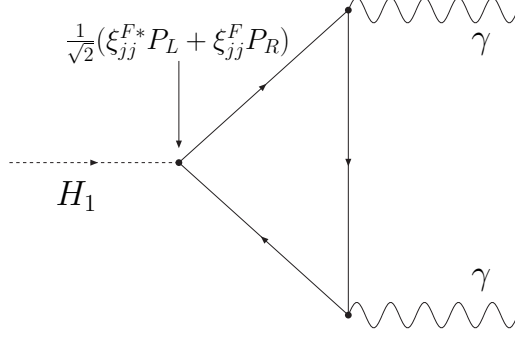


FIG. 1: $H_1 \rightarrow \gamma\gamma$ at one-loop level. The internal lines represent either leptons or quarks.

One sees that H_1, H_2 and h are mass eigenstates, i.e., there is no mixing between the Higgs bosons. To be definite, H_1 is always assumed to be lighter than H_2 , without losing any generality. Now we will see how to make H_1 decaying into both e^+e^- and $\gamma\gamma$ channels with comparable and significant rates.

It is clear from (1) that H_1 decays into e^+e^- at the tree level. By ignoring the electron mass, the decay width is

$$\Gamma(H_1 \rightarrow e^+e^-) = \frac{M_{H_1}}{16\pi} \xi_{ee}^E \xi_{ee}^{E*}, \quad (6)$$

where M_{H_1} is the mass of H_1 and ξ_{ee}^E is the Yukawa coupling between H_1 and electron. Meanwhile, H_1 decays to $\gamma\gamma$ through triangle fermion loop (see Fig. (1)). Notice that the triangle gauge boson loops, which are dominant in the SM, are absent here. This is because that the term $(D_\mu \phi_2)^\dagger (D_\mu \phi_2)$ only results in a four-point vertex $\phi_2^+ \phi_2 B^\mu B_\mu$ due to the vanishing vev of ϕ_2 . Similarly the scalar loops do not appear here either. Fig. (1) yields the following decay amplitude,

$$\begin{aligned} M(H_1 \rightarrow \gamma\gamma) &= 2A^+(pqg^{\mu\nu} - p^\nu q^\mu) \varepsilon_\mu^*(p) \varepsilon_\nu^*(q) - 2A^- i\epsilon^{\mu\nu\rho\sigma} p_\rho q_\sigma \varepsilon_\mu^*(p) \varepsilon_\nu^*(q) \\ &= A^+ F_{\mu\nu} F^{\mu\nu} + iA^- F_{\mu\nu} \widetilde{F}^{\mu\nu}, \end{aligned} \quad (7)$$

with $\widetilde{F}^{\mu\nu} = \frac{1}{2}\epsilon^{\mu\nu\rho\sigma} F_{\rho\sigma}$. There are both CP even amplitude A^+ and CP odd amplitude A^- , as the Yukawa couplings are complex,

$$\begin{aligned} A^+ &= \sum_j \frac{-ie^2 Q_j^2 M_j}{8\pi^2 M_{H_1}^2} \left(\frac{\xi_{jj}^F}{\sqrt{2}} + \frac{\xi_{jj}^{F*}}{\sqrt{2}} \right) \int_0^1 dx \int_0^{1-x} dy \frac{1-4xy}{t_j - xy} \\ A^- &= \sum_j \frac{-ie^2 Q_j^2 M_j}{8\pi^2 M_{H_1}^2} \left(\frac{\xi_{jj}^F}{\sqrt{2}} - \frac{\xi_{jj}^{F*}}{\sqrt{2}} \right) \int_0^1 dx \int_0^{1-x} dy \frac{1}{t_j - xy} \end{aligned} \quad (8)$$

where $t_j = M_j^2/M_{H_1}^2$, $F = E, U$ or D for the virtual fermions to be charged leptons, up quarks or down quarks. Q_j are their charges. The decay width of $H_1 \rightarrow \gamma\gamma$ is thus,

$$\Gamma(H_1 \rightarrow \gamma\gamma) = \frac{M_{H_1}^3}{16\pi}(|A^+|^2 + |A^-|^2) \quad (9)$$

which agree with results in [5].

Now we are in a position to make H_1 decays into e^+e^- and $\gamma\gamma$ with comparable and significant rate. To enhance the loop-suppressed $\gamma\gamma$ channel, the Yukawa couplings of H_1 to heavy leptons and/or quarks should be much larger than that of H_1 to electron. Let us first consider case of a quark-phobic ϕ_2 . We need then a heavy lepton to have a large coupling with ϕ_2 , which we may, for example, choose it to be the τ lepton. Taking the mass of H_1 to be a few hundred GeV, $|\xi_{\tau\tau}^E|/|\xi_{ee}^E| > 10^3$ is required to make the branching ratios for e^+e^- and $\gamma\gamma$ comparable. However with this parameter set, the e^+e^- channel would be too rare to be observable:

$$\text{Br}(H_1 \rightarrow e^+e^-) < \frac{\Gamma(H_1 \rightarrow e^+e^-)}{\Gamma(H_1 \rightarrow \tau^+\tau^-)} \sim \frac{|\xi_{ee}^E|^2}{|\xi_{\tau\tau}^E|^2} < 10^{-6} \quad (10)$$

So a quark-phobic ϕ_2 will not work.

At first glance, the same would happen if ϕ_2 has large couplings to quarks. Naively, $\text{Br}(H_1 \rightarrow e^+e^-) \ll \text{Br}(H_1 \rightarrow q\bar{q})$ would make e^+e^- final state too scarce to be observed. Fortunately, if the mass of H_1 is smaller than 340 GeV, roughly twice of top quark mass, H_1 cannot decay to top quark pairs. In this case, one can make the e^+e^- and $\gamma\gamma$ decay channels comparable and observable, by having $\xi_{tt}^U \gg \xi_{ee}^E$ and keeping the Yukawa couplings between ϕ_2 and the other quarks small. H_1 cannot be too light either, otherwise it should have been observed at LEP already, as the coupling between H_1 and electrons is not too small in this model. For illustrations, we will take (somewhat *ad hoc*) $M_{H_1} = 200$ GeV and $M_{H_2} = M_{H^+} = 300$ GeV in the following numerical discussions.

To get $\Gamma(H_1 \rightarrow e^+e^-)/\Gamma(H_1 \rightarrow \gamma\gamma) \simeq 1$, one needs the ratio of Yukawa couplings to be $\xi_{tt}^U/\xi_{ee}^E \simeq 1800$. To make e^+e^- and $\gamma\gamma$ decay channels observable, we may assume universal coupling of H_1 to all fermions except the top quark². With this parameter set, the main decay channel will be $H_1 \rightarrow gg$. It is then easy to estimate the branching ratios:

$$\text{Br}(H_1 \rightarrow e^+e^-) \simeq \text{Br}(H_1 \rightarrow \gamma\gamma) \sim \frac{\Gamma(H_1 \rightarrow \gamma\gamma)}{\Gamma(H_1 \rightarrow gg)} \sim \frac{\alpha^2}{2\alpha_s^2} \simeq 1/200, \quad (11)$$

² The Yukawa couplings between ϕ_2 and other fermions need only to be of the same order of ξ_{ee}^E . A universal coupling is not necessary, though it makes numerical analysis simpler.

which is not too small and should be measurable. Therefore with such a specific parameter set, it would be possible for Type III 2HDM to produce significant and comparable branching ratios for e^+e^- and $\gamma\gamma$ decays.

In this parameter set, the production of H_1 at hadron collider is mainly through gluon fusion, which is similar to that of the SM Higgs. Therefore, to copiously produce H_1 at an observable rate, one would expect ξ_{tt}^U to be of order one. Generically, one needs ξ_{ee}^E to be no less than 10^{-4} . In the following, we will discuss whether this is possible concerning phenomenological constraints from low energy processes.

III. PHENOMENOLOGICAL CONSTRAINTS AND IMPLICATIONS

With extra not-so-heavy Higgs particles, the type III 2HDM has rich phenomenological implications at low energy. For example, it may contribute to the anomalous magnetic moments of leptons, which have been measured to very high precision. Complex phases of the Yukawa couplings of H_1 to leptons may contribute to leptonic EDMs at one-loop level. Finally, non-diagonal Yukawa couplings of H_1 to leptons would lead to LFV processes, such as $\mu \rightarrow e\gamma$ and $\mu^- \rightarrow e^-e^+e^-$. Non-diagonal Yukawa couplings of H_1 to quarks might produce large hadronic FCNC processes and complex phases may induce CP violating effects in hadrons, which are both tightly constrained experimentally. However these couplings do not affect e^+e^- and $\gamma\gamma$ signals at all, therefore we will simply assume that their effects can be tuned away and not consider them further.

The complex Yukawa matrix for leptonic couplings ξ_{ij}^E contains 18 real parameters, which are obviously too many. So we will consider two extreme scenarios, just for illustration, in numerical discussions:

Scenario I (S1): the leptonic Yukawa matrix is diagonal with a universal CP phase $\pi/6$, i.e. $\xi_{ij}^E = |\xi_{ij}|e^{i\theta_{ij}}$, with $|\xi_{ij}| = 0$ (for $i \neq j$) and $\theta_{ii} = \pi/6$ (for $i = e, \mu, \tau$).

Scenario II (S2) all leptonic Yukawa couplings are assumed to have universal magnitude with CP phase $\pi/6$, namely $\xi_{ij}^E = \xi e^{i\pi/6}$.

Actually, numerical analyses in these two scenarios yield rather reasonable estimates on constraints on the general Yukawa matrix.

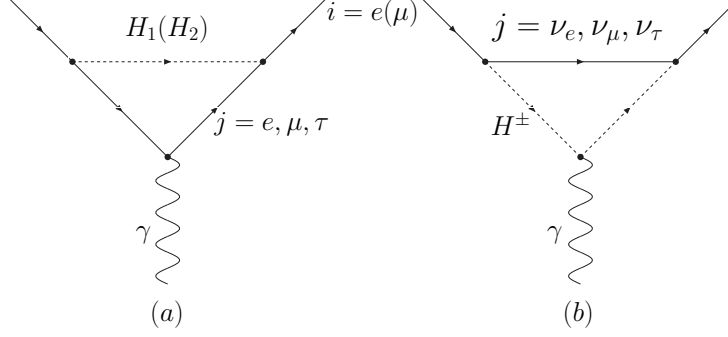


FIG. 2: The one-loop diagrams contribute to $g-2$ and EDM in type III 2HDM.

A. $g - 2$

For the electron, the experimental data gives $a_e^{exp} \equiv (g_e - 2)/2 = (1159652181.1 \pm 0.7) \times 10^{-12}$ [6] which agrees well with the SM calculation $a_e^{SM} = (1159652182.8 \pm 7.7) \times 10^{-12}$ [7]. The extra scalars in 2HDM may contribute to the anomalous magnetic moment at one-loop level, as shown in Fig. (2). We will assume $a_e^{2HDM} < 10 \times 10^{-12}$ which is slightly larger than the theoretical error, to get a restriction on the corresponding Yukawa couplings. For muon, there is some discrepancy between experiments and the SM calculation [6]: $a_\mu^{exp} - a_\mu^{SM} = (292 \pm 63 \pm 58) \times 10^{-11}$, so we assume an upper limit $a_\mu^{2HDM} < 292 \times 10^{-11}$.

Fig. (2) yields the one-loop contribution to $g - 2$ via neutral and charged Higgs as

$$\begin{aligned}
 a_i = & \sum_{j=e,\mu,\tau} \left[\frac{1}{16\pi^2} \int_0^1 dz \frac{-\xi_{ji}^E \xi_{ji}^{E*} M_i^2 z(1-z)}{M_{H^+}^2 - zM_i^2 - i\epsilon} \right. \\
 & + \frac{1}{32\pi^2} \int_0^1 dz \frac{M_i^2 z^2(1-z)(\xi_{ij}^E \xi_{ij}^{E*} + \xi_{ji}^{E*} \xi_{ji}^E) + M_i M_j z^2(\xi_{ji}^{E*} \xi_{ij}^{E*} + \xi_{ij}^E \xi_{ji}^E)}{(1-z)M_{H_1}^2 + zM_j^2 - z(1-z)M_i^2 - i\epsilon} \\
 & \left. + \frac{1}{32\pi^2} \int_0^1 dz \frac{M_i^2 z^2(1-z)(\xi_{ij}^E \xi_{ij}^{E*} + \xi_{ji}^{E*} \xi_{ji}^E) - M_i M_j z^2(\xi_{ji}^{E*} \xi_{ij}^{E*} + \xi_{ij}^E \xi_{ji}^E)}{(1-z)M_{H_2}^2 + zM_j^2 - z(1-z)M_i^2 - i\epsilon} \right] \quad (12)
 \end{aligned}$$

with $i = e, \mu$.

In S1, the expression (12) is proportional to M_e^2 (M_μ^2) for electron (muon) anomalous magnetic moment. That is why, though the upper limit on a_μ^{2HDM} is about 300 times larger than that of a_e^{2HDM} , the constraint $\xi_{\mu\mu} < 0.64$ from $a_\mu^{2HDM} < 292 \times 10^{-11}$ is more severe than the limit $\xi_{ee} < 5.9$ from $a_e^{2HDM} < 10 \times 10^{-12}$.

In S2, the expression (12) is dominated by $M_e M_\tau$ ($M_\mu M_\tau$) term for the electron (muon) anomalous magnetic moment. In this scenario, the electron $g - 2$ leads to $\xi < 0.18$, which

is slightly smaller than the restriction $\xi < 0.21$ from the muon $g - 2$.

B. EDM

The electron and muon EDMs are extremely small in the SM, which arise at three-loop level. But in the present model, the complex Yukawa couplings give rise to non-zero EDMs at one loop level, as shown in Fig.(2). Thus these couplings might be severely constrained by the experimental data. The current experimental measurements cited by PDG2008 [6] are $d_e^{exp} = (0.07 \pm 0.07) \times 10^{-26}$ e cm and $d_\mu^{exp} = (3.7 \pm 3.4) \times 10^{-19}$ e cm. The corresponding upper limits can be found in the original experimental literature: $d_e < 0.16 \times 10^{-26}$ e cm (90% C.L.)[8] and $d_\mu < 10 \times 10^{-19}$ e cm (95% C.L.)[9]. We will use these upper limits, instead of the PDG numbers, to obtain constraints. The leptonic EDM d_i ($i = e, \mu$) has contributions from neutral Higgs bosons H_1, H_2 and charged one H^+ . The contribution from the internal H^+ is suppressed by the neutrino masses, which can be safely dropped. The contribution from the neutral Higgs bosons is

$$d_i = \sum_{j=e,\mu,\tau} \frac{ie}{64\pi^2} \int_0^1 dz \left[\frac{M_j z^2 (\xi_{ij}^E \xi_{ji}^E - \xi_{ji}^{E*} \xi_{ij}^{E*})}{(1-z)M_{H_1}^2 + zM_j^2 - z(1-z)M_i^2 - i\epsilon} - \frac{M_j z^2 (\xi_{ij}^E \xi_{ji}^E - \xi_{ji}^{E*} \xi_{ij}^{E*})}{(1-z)M_{H_2}^2 + zM_j^2 - z(1-z)M_i^2 - i\epsilon} \right], \quad (13)$$

As expected, d_i vanishes when the phase of ξ_{ij}^E is zero and will not yield any constraints on the Yukawa couplings. If only virtual τ lepton is considered inside the loop and the external lepton mass is neglected, our formula agrees with [11]. The above equation is also consistent with [10].

In S1, the expression (13) is proportional to M_e (M_μ) for electron (muon) EDM. It is then straightforward to find $\xi_{ee} < 0.013$ and $\xi_{\mu\mu} < 31$ from the experimental upper limits.

In S2, the expression (13) is dominated by M_τ term for both d_e and d_μ . Therefore the constraint $\xi < 3.9 \times 10^{-4}$ from d_e is significantly tighter than the limit $\xi < 9.6$ from d_μ .

C. LFV processes: $\mu^- \rightarrow e^- \gamma$ and $e^- e^+ e^-$

Both neutral scalars $H_{1,2}$ and charged one H^+ can contribute to $\mu \rightarrow e \gamma$ decay through loop diagrams, as shown in Fig. (3). This channel has not been observed, but with an upper

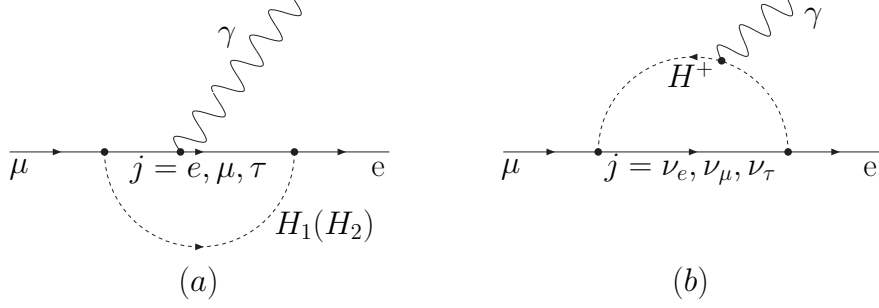


FIG. 3: $\mu \rightarrow e\gamma$ decay in 2HDM at one-loop level.

limit $\text{Br}(\mu^- \rightarrow e^- \gamma) < 1.2 \times 10^{-11}$ [6]. For muon decays into $e^- e^+ e^-$ final state, it may even occur at tree level in our model. Experimentally, the upper limit $\text{Br}(\mu^- \rightarrow e^- e^+ e^-) < 1.0 \times 10^{-12}$ [6] is also very small.

The decay amplitude of $\mu \rightarrow e\gamma$ can be expressed as

$$M(\mu \rightarrow e + \gamma) = \frac{e\varepsilon_\mu^*(q)}{16\pi^2} (-i) \bar{U}_e(p')(p + p')^\mu (A_L P_L + A_R P_R) U_\mu(p), \quad (14)$$

where A_L and A_R are,

$$\begin{aligned} A_L = & \sum_{j=e,\mu,\tau} \int_0^1 dx \int_0^1 dy \int_0^1 dz \delta(x + y + z - 1) \\ & \left[\frac{xzM_\mu \xi_{je}^{E*} \xi_{j\mu}^E}{(1-z)M_{H^+}^2 - xzM_\mu^2 - yzM_e^2} + \frac{-\frac{1}{2}yzM_\mu \xi_{ej}^E \xi_{\mu j}^{E*} - \frac{1}{2}xzM_e \xi_{je}^{E*} \xi_{j\mu}^E - \frac{1}{2}M_j(1-z)\xi_{je}^{E*} \xi_{\mu j}^{E*}}{zM_{H_1}^2 + (1-z)M_j^2 - yzM_\mu^2 - xzM_e^2} \right. \\ & \left. + \frac{-\frac{1}{2}yzM_\mu \xi_{ej}^E \xi_{\mu j}^{E*} - \frac{1}{2}xzM_e \xi_{je}^{E*} \xi_{j\mu}^E + \frac{1}{2}M_j(1-z)\xi_{je}^{E*} \xi_{\mu j}^{E*}}{zM_{H_2}^2 + (1-z)M_j^2 - yzM_\mu^2 - xzM_e^2} \right], \\ A_R = & \sum_{j=e,\mu,\tau} \int_0^1 dx \int_0^1 dy \int_0^1 dz \delta(x + y + z - 1) \\ & \left[\frac{yzM_e \xi_{je}^{E*} \xi_{j\mu}^E}{(1-z)M_{H^+}^2 - xzM_\mu^2 - yzM_e^2} + \frac{-\frac{1}{2}yzM_\mu \xi_{ej}^E \xi_{je}^{E*} - \frac{1}{2}xzM_e \xi_{\mu j}^{E*} \xi_{ej}^E - \frac{1}{2}M_j(1-z)\xi_{ej}^E \xi_{j\mu}^E}{zM_{H_1}^2 + (1-z)M_j^2 - yzM_\mu^2 - xzM_e^2} \right. \\ & \left. + \frac{-\frac{1}{2}yzM_\mu \xi_{ej}^E \xi_{je}^{E*} - \frac{1}{2}xzM_e \xi_{\mu j}^{E*} \xi_{ej}^E + \frac{1}{2}M_j(1-z)\xi_{ej}^E \xi_{j\mu}^E}{zM_{H_2}^2 + (1-z)M_j^2 - yzM_\mu^2 - xzM_e^2} \right]. \quad (15) \end{aligned}$$

In terms of $A_{L,R}$, the decay width can be expressed as,

$$\Gamma(\mu \rightarrow e + \gamma) = \frac{\alpha M_\mu^3}{1024\pi^4} (|A_L|^2 + |A_R|^2). \quad (16)$$

For $e^-e^+e^-$ channel, ignoring the mass of electron, the decay width can be calculated as

$$\Gamma = \frac{|\xi_{ee}^E|^2 \xi_{\mu e}^{E*} \xi_{e\mu}^E M_\mu^5}{24576\pi^3} \left(\frac{1}{M_{H_1}^4} + \frac{1}{M_{H_2}^4} \right). \quad (17)$$

In S1, these LFV processes are forbidden because of the diagonal lepton Yukawa coupling. In S2, we find numerically the restriction $\xi < 1.2 \times 10^{-3}$ from the $e\gamma$ channel, which is smaller than the constraint $\xi < 2.2 \times 10^{-3}$ from the $e^-e^+e^-$ channel.

IV. SUMMARY

We have so far found that, type III 2HDM may produce comparably observable e^+e^- and $\gamma\gamma$ signals simultaneously from decays of neutral scalars. This provides an interesting alternative to the RS model. To make this happen, certain amount of fine tunings are necessary. First, to enhance the loop suppressed $\gamma\gamma$ channel to be comparable with e^+e^- channel, the Yukawa coupling of H_1 to top quark (ξ_{tt}^U) must be a few thousand times larger than those to the rest of fermions. In addition, the mass of H_1 should be smaller than twice of the top quark mass, to prevent H_1 decays predominantly to top-anti-top pair. Under these conditions, H_1 would decay mainly to two gluons, with the branching ratios of e^+e^- and $\gamma\gamma$ channels to be about half a percent.

Notice that these requirements determine only the relative strength of different Yukawa couplings, while their absolute sizes would determine the number of events produced in colliders. If the parameters in these models do meet the above requirements, H_1 will be mainly produced via gluon fusion in hadron colliders, analogous to the SM Higgs. For H_1 to be produced roughly at the same rate as the SM Higgs, ξ_{tt}^U needs to be of order one, which means that the electron Yukawa coupling ξ_{ee}^E should be around $\text{few} \times 10^{-4}$.

Generically, the non-diagonal and complex Yukawa couplings lead to FCNC and CP violating processes. We have thus examined constraints from the anomalous magnetic moment and EDM of leptons, LFV processes $\mu^- \rightarrow e^- \gamma$ and $e^- e^+ e^-$. Because of the large number (18 in total) of free parameters in Yukawa interactions, we have chosen two scenarios as illustration for numerical discussions. In the scenario of diagonal Yukawa couplings with an universal CP phase, we find the most stringent constraint $\xi_{ee}^E < 0.013$ from electron EDM, and $\xi_{\mu\mu}^E < 0.64$ from muon $g-2$ measurement. In the scenario of universal Yukawa couplings, the most stringent constraint $\xi < 3.9 \times 10^{-4}$ comes again from the electron EDM. In both

scenarios, the constraint on ξ_{ee}^E is above the requirement from the dielectron and diphoton signals, so the model survives existing phenomenological tests. The same conclusion also holds true for general Yukawa interactions.

So, if resonances below $2M_t$ are found in experiments to decay into e^+e^- and $\gamma\gamma$ with comparable rates, more work would be needed to delineate their origins. Under this circumstance, we have just seen that type III 2HDM provides one alternative possibility, in addition to the RS-model. It would be interesting to see whether there exist other possibilities.

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